QCD deconfinement – phase transitions and collapsing quark stars.

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Abstract

In this paper we discuss the QCD phase-transitions in the nontopological soliton model of quark confinement and explore possible astrophysical consequences. Our key idea is to look at quark stars (which are believed to exist since the quark matter is energetically preferred over the ordinary matter) from the point of view of soliton model.

We propose that the phase transition taking place during the core collapse of massive evolved star may provide a new physical effect not taken into account in modeling the supernova explosions. We also point out the possibility that merging quark stars may produce gamma-ray bursts energetic enough to be at cosmological distances. Our idea based on the finite-temperature nontopologiocal soliton model overcomes major difficulties present in neutron star merger scenario — the baryon loading problem and nonthermal spectra of the bursts.

1 Introduction.

The possibility that quark matter could be the ground state of hadronic matter attracted a lot of attention since Witten's seminal paper [1].

There are generally two kinds of natural settings in which to seek for macroscopic quark configurations. First of them stems from cosmological quark-hadron phase-transition which occured when the temperature of the early Universe dropped below 100–200 MeV [2]. The second one is of an astrophysical nature and places quark matter inside compact objects like strange stars, cores of neutron stars etc. [3]. These possibilities are by now rather well studied (see [4] for a review). The basic physical model adopted in studies of the quark matter is the MIT bag model for hadrons. In this picture the quark star (or primordial quark lump) is understood as an enormous "hadron" i.e. an ensamble of quarks inside a confining bag. The notion of bag is believed to capture, at a phenomenological level the essence of (yet unknown precisely) zero-tempetrature nonperturbative QCD effects. The simplest way to generalize the MIT model to a finite temperature theory is to adopt the Friedberg-Lee nontopological soliton model [5]. In this picture the "bag constant" is again a difference between the values of an effective potential U_{eff} in perturbative QCD vacuum and in physical vacuum respectively. However the effecitve potential is temperature-dependent now and so are its extrema. Hence the bag constant is no longer a constant but rather a function of temperature T (and chemical potential μ). This opens a possibility of phase-transitions in macroscopic quark configurations. Phase-transitions in a nontopological soliton model of hadrons has been studied by Song,Enke and Jian Ong [6] with hope to understand the confinement. Our investigation performed in the next section is in a similar vein. We demonstrate that macroscopic quark configurations may undergo phase transitions at some critical temperature releasing thereby a huge amount of energy. Then in section 3 we contemplate possible astrophysical consequences of such phase-transitions. Section 4 contains concluding remarks.

2 Phase transitions in Friedberg-Lee quark solitons.

We focus our attention on the Friedberg-Lee nontopological soliton model defined by the Lagrangian density [5]

$$\mathcal{L} = \frac{1}{2}\partial_\mu\varphi\partial^\mu\varphi - U(\varphi) + \overline{\psi}(i\gamma^\mu D_\mu - m - g_0\varphi)\psi$$

where ψ is the quark - field, m – the quark mass and φ is a phenomenological scalar field believed to represent nonperturbative QCD effects. The self-interaction potential of the scalar field φ is assumed to be of the form:

$$U(\varphi) = \frac{a}{2!}\varphi^2 - \frac{b}{3!}\varphi^3 + \frac{c}{4!}\varphi^4 + B. \tag{2.1}$$

The constants a, b, c and g_0 represent four parameters of the model. The potential $U(\varphi)$ has typically two minima, one at $\varphi_0 = 0$ which is a local minimum associated with the perturbative vacuum state and the second — absolute minimum, at

$$\varphi_v = \frac{3|b|}{2c}\left[1 + (1 - \frac{8ac}{b^2})^{\frac{1}{2}}\right] \tag{2.2}$$

corresponding to the physical vacuum. The constant B measures the difference in potential between two vacuum states and thus has the same meaning as the MIT bag constant. However unlike in the MIT model the confining bag is not put in á priori but rather results from the dynamics of the scalar field φ.

Due to nonlinearity of the potential $U(\varphi)$ the Lagrangian (2.1) leads to nontopological soliton solutions. There are two limiting types of the shape of $U(\varphi)$ for which the soliton solution exists. The first represents degenerate vacuum state with both minima at the same level $U(\varphi_v) = U(\varphi_0)$ and consequently $B = 0$. The second case is the one when the local minimum and maximum overlap producing thus an inflection point. This is the moment when the perturbative and physical vacuum states become metastable. Soliton solutions exist for all intermediate potential shapes i.e. for $U(\varphi)$ with two minima and one maximum. At this stage the demand that soliton solution exist imposes a restriction on the a, b, c – parameters of the model.

It is now interesting to study how the properties of soliton solutions change with temperature. For this purpose it is convenient to perform the decomposition of the scalar field φ into thermodynamical average

$$\sigma = <\varphi> = \frac{Tre^{-\beta H}\varphi}{Tre^{-\beta H}}$$

and quantum corrections $\tilde{\varphi}$:

$$\varphi = \sigma + \tilde{\varphi}. \tag{2.3}$$

Consequently the potential function $U(\varphi)$ transforms into an effective potential $U_{eff}^{\beta}(\sigma)$ which is explicitly temperature-dependent:

$$U_{eff}^{\beta}(\sigma) = U_{eff}^{T=0}(\sigma) + \delta_{\beta}U_{eff}(\sigma) \tag{2.4}$$

where $U_{eff}^{T=0}(\sigma)$ is the zero-temperature one-loop effective potential and the one-loop temperature dependent quantum correction reads:

$$
\begin{aligned}
\delta_{\beta}U_{eff}(\sigma) &= \frac{1}{\beta}\int \frac{d^3k}{(2\pi)^3}ln(1 - e^{-\beta E_M}) \\
&= \frac{1}{2\pi^2\beta^4}\int_0^{\infty} x^2 dx ln(1 - e^{-(x^2+\beta^2 M^2)^{\frac{1}{2}}})
\end{aligned}
$$

$$\tag{2.5}$$

where $x = \beta k$, k is the wave number, $E_M^2 = k^2 + M^2$ and M is the effective mass of the scalar field i.e. equal to the second derivative of $U_{eff}^{\beta}(\sigma)$ taken at the global minimum.

The expression (2.5) is in a sense incomplete since the contributions coming from fermions have been neglected. However such thermal corrections are supposed to be significant only at temperatures exceeding the Fermi level $kT >> E_F$ and an inspection at Table 2 shows that indeed Fermi energies are of order of several hundreds of MeV which is much greater than thermal energies encountered during the core collapse.

Expanding $\delta_{\beta}U_{eff}(\sigma)$ in power series around $\beta M = 0$ in the integral (2.5) yields the well known expresion

$$\delta_{\beta}U_{eff}(\sigma) = \frac{1}{2\pi^2\beta^4}(\int_0^{\infty} x^2 dx ln(1 - e^{-x})$$

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$$+ \quad \beta^2 M^2 \frac{1}{2} \int_0^\infty dx \frac{x}{(e^x - 1)} + O(\beta^4 M^4)) \tag{2.6}$$

which can be integrated to give the following formula:

$$\delta_\beta U_{eff}(\sigma) = -\frac{\pi^2}{90}T^4 + \frac{T^2}{24}M^2. \tag{2.7}$$

Therefore the effective potential $U_{eff}(\sigma, T)$ is equal to

$$\begin{aligned}
U_{eff}(\sigma, T) &= -\frac{bT^2}{24}\sigma + \frac{1}{2}(a + \frac{cT^2}{24})\sigma^2 - \frac{1}{6}b\sigma^3 \\
&+ \frac{c}{24}\sigma^4 + \frac{a}{24}T^2 + \frac{c}{1152}T^4 - \frac{\pi^2}{90}T^4 + B
\end{aligned} \tag{2.8}$$

The temperature dependence of the effective potential $U_{eff}(\sigma, T)$ is illustrated on Figure 1. For $T = 0$ the effective potential has three extrema corresponding to the perturbative vacuum, local maximum (potential barrier) and physical vacuum, respectively. At some critical temperature T_0 located here somewhere between $0.3\ m_s$ and $0.4\ m_s$, where $m_s = 197.32$ MeV is the strange quark mass, the perturbative vacuum and the local maximum overlap, inflection point develops and bag constant vanishes. This is also the last moment when the soliton exists. Values of the temperature T_0 for each set of parameters are presented in Table 1.

	a(fm^{-2})	b(fm^{-1})	c	T_0(MeV)
1	51.60	799.90	4000	68.71
2	7.671	107.27	500	84.41
3	12.85	196.34	1000	77.28
4	40.88	783.80	5000	60.46
5	66.42	1411.60	10000	55.59
6	107.32	2537.00	20000	49.52
7	321.75	9824.80	100000	38.67

Table 1: Values of the critical temperature T_0 for selected set of parameters.

Alternative way of illustrating this transition is presented in Figure 2 where the location of extrema of the effective potential is plotted against the temperature.

For temperatures greater than T_0 the soliton disappears and the structure of the vacuum changes. By comparing the effective potential (2.8) with the potential (2.1) one may identify last four terms in (2.8) with temperature dependent "bag constant" $B(T)$ describing the pressure of the physical vacuum exerted on the false vacuum

$$B(T) = \frac{a}{24}T^2 + \frac{c}{1152}T^4 - \frac{\pi^2}{90}T^4 + B. \tag{2.9}$$

Fig.3 shows $B(T)$ as a function of temperature calculated for the set of parameters no.1 from Table 1. At the critical temperature $B(T)$ is nonzero (equal to $\sim 0.1 \times m_s^4$ in this

	B	k_F
ad **1**	$(100 MeV)^4$	277 MeV
	$(180 MeV)^4$	499.3 MeV
ad **2**	$(100 MeV)^4$	306.9 MeV
	$(180 MeV)^4$	552.6 MeV

Table 2: Fermi levels for selected model parameters.

specific example) despite the fact that the soliton disappears. It means that the energy contained in the bag soliton is liberated at the critical temperature. This observation is crucial for possible astrophysical scenario envisaged in this paper.

Before closing this section let us discuss the behavior of the total energy of the quark core. The total energy of a soliton bag (including self-gravity) takes the following form

$$E = \frac{\alpha N^{\frac{4}{3}}}{R} + \frac{4}{3}\pi BR^3 - \frac{3}{5}\frac{G}{R}(\frac{\alpha N^{\frac{4}{3}}}{R} + \frac{4}{3}\pi BR^3)^2, \tag{2.10}$$

where $\alpha = \frac{9}{4}(\frac{9\pi}{4})^{1/3}$. Following [7] it is convenient to introduce a dimensionless quantity

$$x = (\frac{4\pi B}{\alpha})^{\frac{1}{4}} N^{\frac{-1}{3}} R \tag{2.11}$$

Then one has,

$$E = \alpha^{\frac{3}{4}}(4B\pi)^{\frac{1}{4}} N(\frac{1}{x} + x^3 - \frac{e}{x}(\frac{3}{x} + x^3)^2) \tag{2.12}$$

with

$$e = \frac{1}{5}GN^{\frac{2}{3}}(4\pi\alpha B)^{\frac{1}{2}}. \tag{2.13}$$

Figure 4 illustrates the dimensionless energy function $f(x)$

$$f(x) = \frac{1}{x} + x^3 - \frac{e}{x}(\frac{3}{x} + x^3)^2 \tag{2.14}$$

represented by a solid curve. It has a local minimum corresponding to the stable core configuration where the repulsive term coming from the Fermi statistics obeyed by the quarks is counterbalanced by the joint action of bag pressure and gravity. We stress that such equilibrium configuration may not be realised during the collapse which is a rapidly progressing dynamical process. One may imagine that while the core is collapsing one is moving along the curve $f(x)$ to the left. When the effective bag constant $B(T)$ vanishes the bag pressure vanishes as well and only much weaker effect of gravity remains to act against fermion repulsion. This changes qualitatively the $f(x)$-function and the new situation is denoted by curve $f_0(x)$ in Figure 3. Hence at the moment when $B(T)$ vanishes the point representing the energetic state of the core jumps (at some x_{cr}) from the curve $f(x)$ to $f_0(x)$ liberating some energy in this process. Moreover the repulsive, fermionic term is dominant now and causes the expansion of the core.

6

3 Possible astrophysical consequences of high-temperature QCD phase–transitions.

The history of science in the present century (or at least its second half) reveals a successful search for the manifestation of fundamental interactions in astrophysical and cosmological context. Consequently for any branch of physics one may find a corresponding branch of astrophysics [8]. With this perspective in mind we outline in this section some astrophysical processes in which the QCD effects discussed above are likely to be operating.

3.1 Core collapse and the supernova explosion.

Massive stars are known to end their evolution with the destabilization and collapse of their iron cores. Following the core collapse a supernova explosion releases most of the matter (and the internal energy) of the progenitor. Although the phenomenon of supernova is known (observationally) since ancient times and despite the tremendous progress achieved in the theory of stellar evolution (especially in the field of numerical simulations) the detailed mechanism of explosion is not well understood.

The most commonly accepted picture is the following: the destabilization of the iron core leads to core collapse during which the core divides into a subsonic, homologously collapsing inner part and outer free-falling part. The collapse of the inner part is halted at nuclear densities by the stiffening of the equation of state and after a rebounce an outwardly propagating shock wave is generated. Numerical simulations show that usually the shock wave is damped mainly due to photodisintegration of Fe-like elements and becomes a standing accretion shock. As a way out of this trouble it has been proposed by Wilson [9] that stagnated shock may be revived by the accretion induced heating due to neutrino absorbtion. Simple criterion for the shock revival can be established in terms of neutrino luminosity L_ν and the average of the rms ν_e energies $\langle \varepsilon_\nu^2 \rangle$ [10]:

$$L_{51}\langle \varepsilon_{15}^2 \rangle > 48 \frac{M_{1.4}^{3/2}}{r_3^{1/2}}, \qquad (3.15)$$

where: $L_{51} = \dfrac{L_\nu}{10^{51} \ erg \ s^{-1}}$, $\varepsilon_{15} = \dfrac{\varepsilon_\nu}{15 \ MeV}$, $M_{1.4} = \dfrac{M}{1.4 \ M_\odot}$ and $r_3 = \dfrac{r}{10^3 \ cm}$. It means that $L_{51}\langle \varepsilon_{15}^2 \rangle$ should be large enough in order to revive the shock. This is not always easy to acheive in a way proposed by Wilson i.e. by neutrino absorbtion. However any other mechanism for the revival is good provided the heating rate exceeds the value

$$\frac{d\varepsilon}{dt} = 53 \frac{M_{1.4}^{3/2}}{r_3^{1/2}} \quad MeV \ s^{-1}. \qquad (3.16)$$

The need for a robust mechanism of explosion is even more pronounced in the context of the Supernova 1987A because the progenitor here was an 18–20 $M_\odot$ blue supergiant and it is commonly accepted that the prompt shock wouldn't work in such massive star [11].

Various phase transitions are expected to occur in the core. Above the nuclear density there may be first or second order phase transitions such like pion and kaon condenstes [11] or deconfined quark matter [12]. There were attempts to include this complicated

physics into numerical simulations of the core collapse but the results are clearly still controversial in the aspect of explosion. Although some progress has recently been achieved in hydrodynamical treatment of neutrino transport it cannot be excluded that physical effects outlined in this paper may find application in realistic models of core collapse.

A proto-neutron star core is a favorable environment for the conversion of ordinary hadronic to strange quark matter through a variety of fluctuations [3]. It seems reasonable to contemplate the properties of the quark matter in the collapsing core within the finite-temperature formalism of Friedberg and Lee. The results of Cooperstein [13] are encouraging in this respect since his simulations demonstrated that the temperature of collapsing core may easily acquire 80 MeV or higher. One may thus expect the occurence of phase transitions discussed in previous section as an additional trigger for supernova explosion. As shown in Fig.3 the energy density that may be liberated at the critical point is of order of $\sim 0.09 \ (m_s)^4 \approx 197.33 \ MeV/fm^3$ and assuming the radius of the quark core of order of $\sim 10 \ km$ this gives the total energy equal to $\sim 10^{53} \ erg$.

3.2 Gamma-bursts from strange stars' mergers.

Since their discovery about 20 years ago [14] gamma ray bursts (GRBs) still remain the most mysterious phenomena in the sky [15]. Recent results provided by the ongoing mission of the Compton GRO [16] demonstrated an isotropic distribution of GRBs on the celestial shpere which combined with their radial distribution being uniform out to a certain distance and falling off beyond strongly supported the idea that GRBs originate from cosmological sources [17]. More recently the correlation between the duration the strength and the hardness of the bursts has been found [18] as predicted by the cosmological scenario. In the light of enormous diversity of durations, time variability and spectra of the GRBs no detailed theoretical model gained general acceptance yet. There is however a number of reasons for believing that the GRBs may come from neutron star mergers. Neutron star binaries such like the famous binary pulsar PSR 1916+13 are known to exist and because they radiate away gravitational waves as predicted by general relativity they will finish their lives in a catastrophic merger event powerful enough to be responsible for GRBs at cosmological distances. The expected rate of neutron star mergers (inferred from the known three binary pulsars) is about $10^{-5.5\pm.5}$ /year/galaxy [19]. This is in perfect agreement with a local rate of $\sim 10^{-6}$ /year/galaxy seen in the BATSE experiment.

The idea that the strange quark matter is the ground state of matter suggests that neutron stars may in fact be hybrid stars consisting of quark central region covered by a layer of ordinary matter [3, 4]. Consequently the binary systems made of quark stars may exist as well. The idea that the GRBs could be explained by the collision of two strange stars has been put foreward by Haensel, Paczyński and Amsterdamski [20]. Their line of thinking was following: a merger of two strange stars thermalizes the kinetic energy of the collision which is roughly about 10^{53} ergs. then the newly born post-merger starnge star radiates most of the thermal energy in the form of neutrinos and antineutrinos within a few seconds after a collision. The neutrino-antineutrino pairs convert into electron-positron pairs which annihilate into gamma rays. Because the huge amount of energy is deposited in a small region (surface of post-merger star) on a time scale of a fraction of a second the e^+, e^- plasma will be thermalized and will produce a fireball with a blackbody spectrum peaking at energies ~ 10 MeV. The fireball model faces two serious problems:

the effect of baryons and the origin of the observed nonthermal spectrum of GRBs. The baryon loading resulting in optically thick fireball and subsequent conversion of radiation into kinetic energy of the fireball is naturally overcome in the case of quark stars (no baryons) but the second objection remains.

Let us consider the merger of two quark stars from the point of view of nontopological soliton approach presented in the section 2. Assume that they are initially cool stable configurations which masses and radii are determined by the bag constant B : [7]

$$
\begin{align}
M_{1,2}(B) &= \mu B^{1/4} N_{1,2} \\
R_{1,2}(B) &= \rho B^{-1/4} N_{1,2}^{1/3}
\end{align}
$$

where $\mu = 4(\frac{3}{2})^{2/3}\pi^{1/2}$, N is the total number of fermions (quarks) and $\rho = \frac{1}{8}(\frac{3}{2})^{5/3}\pi^{1/12}$.

For the post-merger object one should not take zero-temperature bag constant B but rather the effective bag constant $B(T_{pm})$ where the post-merger temperature T_{pm} is of order of 10^{11} K [20] since the kinetic energy of the collision is thermalized. Therefore the mass of the final configuration $M_{pm} = \mu B(T_{pm})^{1/4}(N_1+N_2)$ is lower than the sum of initial masses $M_1 + M_2$ and this mass deficit $\Delta M_{pm} = \mu[B(T_{pm})^{1/4} - B(T=0)^{1/4}](N_1 + N_2)$ is radiated away resembling somewhat the process of nuclear fussion. The phenomenological scalar field σ undergoes out of equilibrium decay into radiation via a one-loop-fermion process $\sigma \to 2\gamma$ [21]. Moreover it is oversimplified to assume that the post-merger configuration is stable against collapse. This opens the possibility that the phase-transitions described in the Section 2 may be operating in this scenario as well. It should be stressed that the second-order phase transition which is seen in the behavior of the effective bag constant $B(T)$ is closely related to the behavior of the phenomenological scalar field σ. At the critical point the scalar field σ (or rather its excess over the hadronic matter case) should be radiated away in a process $\sigma \to 2\gamma$. This radiation would be nonthermal and could explain the nonthermal spectra of the GRBs which are otherwise hard to explain. The detailed calculations of spectra in the merging quark stars scenario will be presented in another paper.

4   Conclusions

If one takes the idea of quark matter seriously then it is natural to consider the collapsing core of a massive evolved star as a prime site for the quark matter formation. Because the temperatures typical for this environment are high one concludes that the finite-temperature theory is more appropriate than zero-temperature MIT model widely used in the literature. Consequently one may expect that massive compact objects produced in the final stages of stellar evolution may contain quark matter (the so called hybrid stars) or be composed exclusively out of it. If one adopts the philosophy that formation of degenerate gas provides the main force stabilizing the remnants of stellar evolution against self-gravity, then we acquire a new link in the chain of stable configurations: white dwarfs (degenerate electrons) $\to$ neutron stars (degenerate neuron gas) $\to$ quark stars (degenerate quark gas) $\to$ black hole .

In this paper we tried to look at the problem of macroscopic quark configurations from the point of view of finite-temperature theory — the so called Friedberg-Lee soliton model. We have demonstrated the possibility of a phase transition during which the confining

bag disappears thus liberating a huge amount of energy. This may occur at temperatures of order of $\sim 40 - 80 \ MeV \approx 5 - 9 \times 10^{11} \ K$ which is a range accessible during the core collapse. The aforementioned phase transition may provide an important physical effect (additional source of energy?) during the supernova explosion. We have also presented an idea according to which merging quark stars discussed in the finite-temperature theory may explain the gamma-ray bursts. Such a model would easily overcome main difficulties that plagued conventional models such like baryon loading and production of nonthermal radiation out of a thermalized fireball.

Acknowledgements

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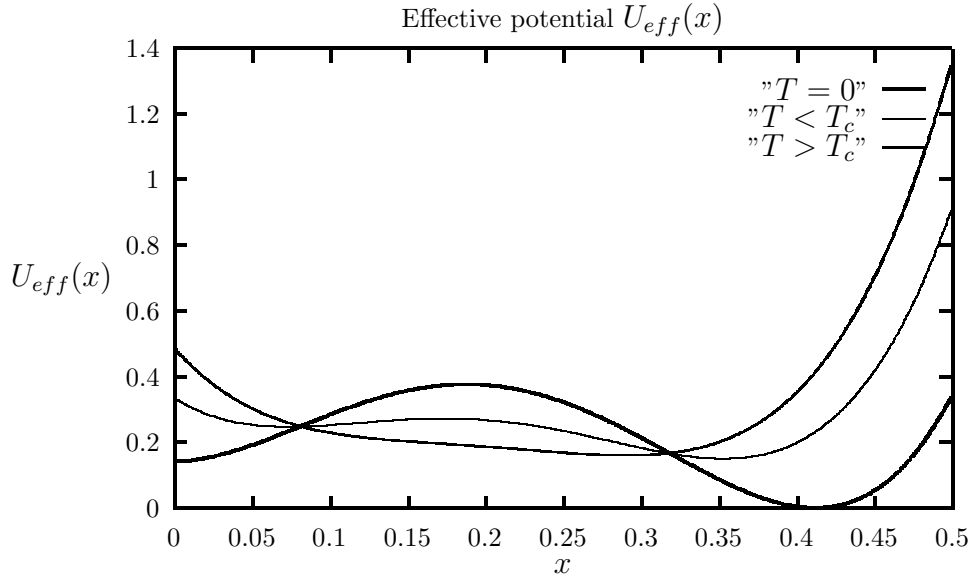
Figure captions:

Figure 1.
Temperature dependence of the effective potential $U_{eff}(x,t)$, ($t = \dfrac{T}{s_0}$ and $x = \dfrac{\sigma}{s_0}$ where $s_0 = 197.32$ MeV is the strange quark mass.)

Figure 2.
Values of σ at the extremum of the effective potential plotted vs. temperature – the qualitative change from three to one extremum occurs at critical temperature.

Figure 3.
The temperature dependence of the bag constant $B(t)$

Figure 4.
The effective energy of the bag loaded with quarks $f(x)$, $f_b(x)$ - the energy of the empty bag, $f_o(x)$ - the energy of free quarks.

Figure 1: Temperature dependence of the effective potential $U_{eff}(x, t)$, ($t = \dfrac{T}{s_0}$ and $x = \dfrac{\sigma}{s_0}$ where $s_0 = 197.32$ MeV is the strange quark mass.)
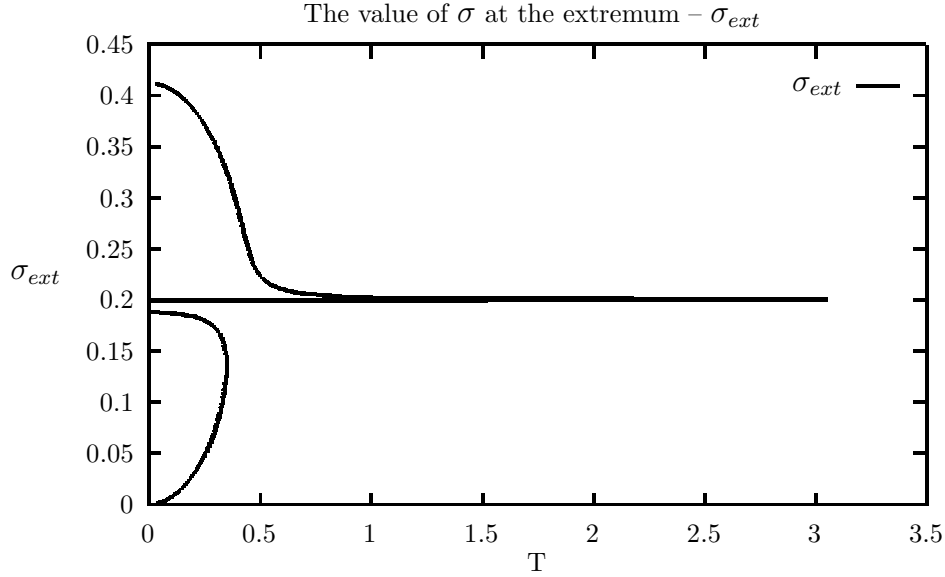

Figure 2: Values of σ at the extremum of the effective potential plotted vs. temperature – the qualitative change from three to one extremum occurs at critical temperature.
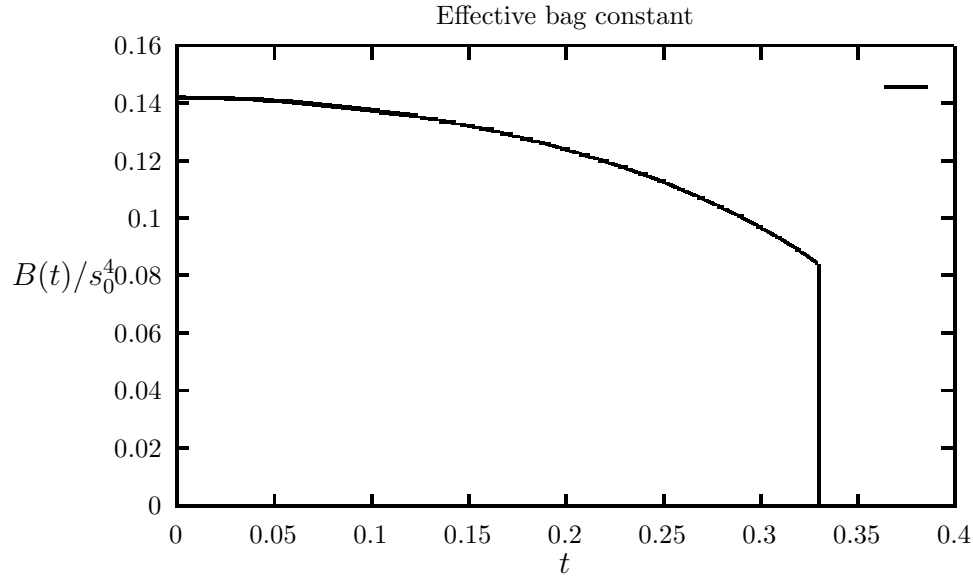
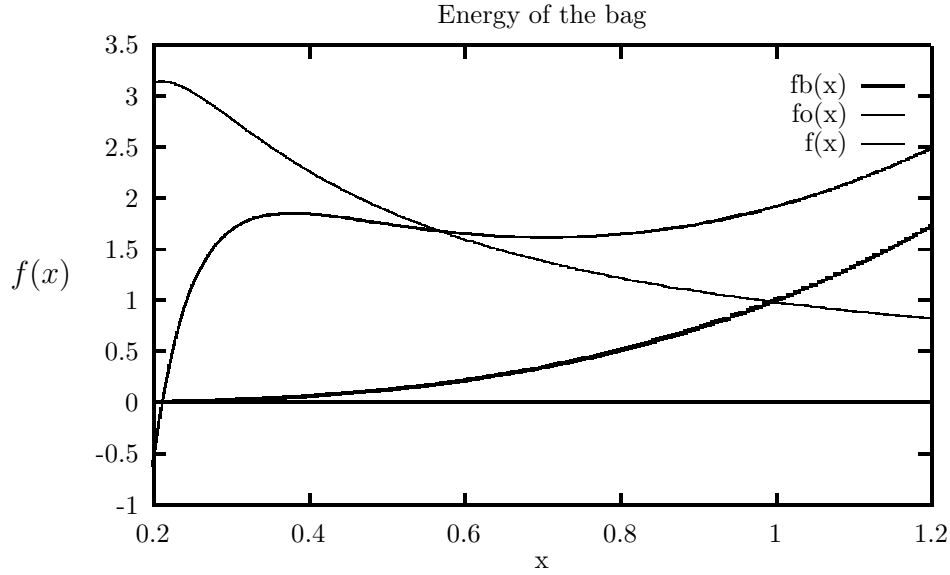Figure 3: The temperature dependence of the bag constant $B(t)$

Figure 4: The effective energy of the bag loaded with quarks $f(x)$, $f_b(x)$ - the energy of the empty bag, $f_o(x)$ - the energy of free quarks.